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$D^0 - \bar{D}^0$ Mixing in the Presence of Isosinglet Quarks

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Abstract

We analyse $\Delta C = 2$ transitions in the framework of a minimal extension of the Standard Model where either a $Q = 2/3$ or a $Q = -1/3$ isosinglet quark is added to the standard quark spectrum. In the case of a $Q = 2/3$ isosinglet quark, it is shown that there is a significant region of parameter space where $D^0 - \bar{D}^0$ mixing is sufficiently enhanced to be observed at the next round of experiments. On the contrary, in the case of a $Q = -1/3$ isosinglet quark, it is pointed out that obtaining a substantial enhancement of $D^0 - \bar{D}^0$ mixing, while complying with the experimental constraints on rare kaon decays, requires a contrived choice of parameters.

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1 Introduction

Flavour-changing processes leading to $K^0 - \bar{K}^0$ and $B^0 - \bar{B}^0$ mixings have played an important rôle in testing the Standard Model (SM) and in constraining some of the physics beyond the SM. Indeed the $K_L - K_S$ mass difference was used to predict the approximate value of the charm quark mass [1] while the experimental value of $B_d - \bar{B}_d$ mixing [2] provided the first hint that the top quark is much heavier than anticipated. Regarding $D^0 - \bar{D}^0$ mixing, its main interest stems from the fact that it provides a sensitive probe of physics beyond the SM. Within the SM the short-distance contributions to $D^0 - \bar{D}^0$ mixing arise from box-diagrams and are small, $\Delta M_D^{\text{box}}(\text{SM}) \approx 10^{-17}$ GeV. Initially, it was argued that the mass difference in the $D^0 - \bar{D}^0$ sector was mainly a long distance effect and a discussion of both short and long distance effects led to the estimate [3,4] that the long distance contribution is about two orders of magnitude larger than the box diagram contribution. A subsequent study suggested [5] that there are significant cancellations among the dispersive channels, leading to the prediction that the long distance contribution is smaller than previously estimated. This conjecture was recently confirmed by an analysis of $D^0 - \bar{D}^0$ mixing in heavy quark effective field theory (HQEFT), including leading order QCD corrections, which led to the result [6] $(\Delta M_D)_{\text{HQEFT}} \approx (0.9-3.5) \cdot 10^{-17}$ GeV. These results imply that the observation of $D^0 - \bar{D}^0$ mixing at the next round of experiments at a tau/charm factory, at Fermilab and at CESR, would provide a clear indication of physics beyond the SM, since the present experimental bound is still four orders of magnitude above the SM prediction and the expected sensitivity of these experiments would easily allow for an improvement of two orders of magnitude. The simplest mechanism to obtain significant new contributions to $D^0 - \bar{D}^0$ mixing consists of having either scalar or gauge flavour-changing neutral currents (FCNC). The importance of Higgs couplings to FCNC and their impact on $D^0 - \bar{D}^0$ mixing have been recently emphasized by Hall and Weinberg [7]. Models with isosinglet quarks provide the simplest framework where Z flavour-changing neutral currents (ZFCNC) arise at tree level, together with deviations from unitarity of the CKM matrix.

In this paper, we will analyse $D^0 - \bar{D}^0$ mixing within a minimal extension of the SM with either a $Q = 2/3$ or a $Q = -1/3$ isosinglet quark. The description of the model is provided in section 2. In section 3 we analyse in detail the contribution to $D^0 - \bar{D}^0$ mixing and point out that in models with a $Q = 2/3$ isosinglet quark and for plausible values of the parameters, the mixing can be sufficiently enhanced to be observed at the next round of experiments. This is to be contrasted with the case of models with a $Q = -1/3$ isosinglet quark, where we show that achieving a significant enhancement of $D^0 - \bar{D}^0$ mixing while satisfying the experimental constraints on rare kaon processes requires fine-tuning of parameters. We also consider, in section 4, the possibility of having an enhancement of single top production in e^+e^- , having in mind a search for this process in LEP II. However, we show that only for a somewhat contrived choice of parameters, is this search feasible at LEP II. In section 5, we present our conclusions.

2 Model with $Q = 2/3$ Isosinglet Quark

In order to settle the notation and provide a framework to discuss ZFCNC, we briefly describe a minimal model where ZFCNC arise in the up quark sector. We will consider a minimal extension of the SM which consists of adding a charge $(2/3)$ quark T whose left-handed and right-handed components are both singlets under $SU(2)$. One may give appropriate mass to all quarks using only one Higgs doublet as in the SM. In this case, apart from the SM Yukawa couplings one has the $SU(2) \times U(1)$ invariant mass terms $M_j \bar{T}_L u_{R_j}$. A slightly more complicated Higgs structure is required in order to obtain spontaneous CP breaking and a possible solution to the strong CP problem. In that case, it has been shown [8] that the minimal Higgs structure consists of introducing a complex $SU(2) \times U(1)$ singlet, together with the SM Higgs doublet. Our analysis does not depend on whether one introduces only a Higgs doublet or a Higgs doublet and a singlet.

Without loss of generality, one may choose a weak-basis where the 3×3 down-quark mass matrix is diagonal, real and the 4×4 up-quark mass matrix has the form:

$$\mathcal{M}_u = \begin{bmatrix} G & J \\ 0 & M \end{bmatrix} \quad (2.1)$$

where G is a (3×3) matrix, while J is a (3×1) column matrix. M corresponds to a good approximation to the mass of the isosinglet quark T . If we denote by m the mass scale of G, J , then it is natural to assume $M \gg m$, since G, J are $\Delta I = 1/2$ mass terms while M is a $\Delta I = 0$ mass term. The weak-eigenstates (u_i^0, T^0) and mass-eigenstates (u_i, T) are related through the unitary transformation:

$$\begin{pmatrix} u^0 \\ T^0 \end{pmatrix}_L = W \begin{bmatrix} u_i \\ T \end{bmatrix}_L. \quad (2.2)$$

It is convenient to write the 4×4 unitary matrix W as:

$$W = \begin{bmatrix} K & R \\ S & X \end{bmatrix} \quad (2.3)$$

where $(K \ R)^\dagger$ is the (4×3) CKM matrix, while K^\dagger corresponds to its (3×3) block connecting standard quarks.

The charged-current interactions can then be written as:

$$\frac{g}{\sqrt{2}} \left[\bar{u}_L K^\dagger \gamma_\mu d_L + \bar{T}_L R^\dagger \gamma_\mu d_L \right] W^\mu + h.c. \quad (2.4)$$

A nice feature of this class of models is that both deviations from unitarity of the matrix K^\dagger and ZFCNC among standard quarks are naturally suppressed [9] by the ratio m^2/M^2 . Indeed the unitarity of W implies:

$$(V_{CKM}) (V_{CKM}^\dagger) = 1 - S^\dagger S \quad (2.5)$$

where $V_{CKM} \equiv K^\dagger$. An approximate diagonalization of $\mathcal{M}_u \mathcal{M}_u^\dagger$ leads to:

$$S \approx -\frac{J^\dagger K}{M} \quad (2.6)$$

and therefore S is of order m/M .

The neutral current interactions are given by:

$$\mathcal{L}_Z = \frac{g}{2 \cos \theta_W} \left[z_{\alpha\beta} \bar{u}_{\alpha L} \gamma^\mu u_{L\beta} - \delta_{ij} \bar{d}_{iL} \gamma^\mu d_{jL} - 2 \sin^2 \theta_W J_{\text{em}}^\mu \right] Z_\mu \quad (2.7)$$

where $\alpha = 1 \dots 4$, with $u_4 \equiv T$ and $u_i \equiv (u, c, t)$. Of particular interest to us are the ZFCNC connecting standard quarks, which are given by:

$$Z_{ij} = -S_i^* S_j \quad (2.8)$$

where the indices i, j in Z_{ij} refer to u, c, t . A crucial question is: for a given value of M , what are the expected values of Z_{ij} ? In order to answer this question, we write explicitly the expressions for the S_i :

$$\begin{aligned} S_1 &\cong -\frac{1}{M} (J_1^* V_{ud}^* + J_2^* V_{us}^* + J_3^* V_{ub}^*) \\ S_2 &\cong -\frac{1}{M} (J_1^* V_{cd}^* + J_2^* V_{cs}^* + J_3^* V_{cb}^*) \\ S_3 &\cong -\frac{1}{M} (J_1^* V_{td}^* + J_2^* V_{ts}^* + J_3^* V_{tb}^*). \end{aligned} \quad (2.9)$$

It is clear from Eq. (2.9) that the values of the S_i crucially depend on the values of the J_i . One may be tempted to consider that the J_i follow the standard up quark hierarchy. We would like to point out that this is not the case in general, since the J_i are independent parameters, unrelated to either the quark mass spectrum or the (3×3) V_{CKM} matrix. This can be easily shown by noting that to leading order in m^2/M^2 , the following relation holds [9]:

$$K^{-1} G G^\dagger K = \bar{m}^2 \quad (2.10)$$

where $\bar{m}^2 \equiv \text{diag} (m_u^2, m_c^2, m_t^2)$. Therefore, both the spectrum of the standard up quark masses and the (3×3) block of the CKM matrix, ($V_{CKM} \equiv K^\dagger$), are determined, to leading order, by the hermitian matrix $G G^\dagger$. As a result, the J_i should be treated as independent parameters, unrelated to the standard quarks mass spectrum. Thus in the next section where we analyse the size of $D^0 - \bar{D}^0$ mixing, we will consider various hypothesis for the relative size of the J_i .

We have described the main features of a model with a $Q = 2/3$ isosinglet quark. It is clear that entirely analogous considerations apply to a model with a $Q = -1/3$ isosinglet quark.

3 $D^0 - \bar{D}^0$ mixing

Within the SM, the short distance contributions to $D^0 - \bar{D}^0$ mixing arise from box-diagrams and are small due essentially to the fact that the dominant terms are those associated to the intermediate quark lines s and d where the suppression factor $(m_s^2 - m_d^2)^2/(M_W^2 m_c^2)$ appears [10]. This factor should be compared to the corresponding factor for $K^0 - \bar{K}^0$ mixing which is m_c^2/M_W^2 (roughly 1300 times higher for $m_s \simeq 0.2$ GeV and $m_c \simeq 1.2$ GeV).

Next we evaluate $D^0 - \bar{D}^0$ mixing in models with either a $Q = 2/3$ or a $Q = -1/3$ isosinglet quark.

3.1 $D^0 - \bar{D}^0$ mixing in a model with a $Q = 2/3$ isosinglet quark

In this case, there are ZFCNC in the up quark sector which lead to a new contribution to $D^0 - \bar{D}^0$ mixing from Z exchange tree graphs (Fig. 1). Using Eq. (2.7) one obtains for the effective $\Delta C = 2$ Lagrangian:

$$\mathcal{L}_{\text{eff}}^Z = \frac{-g^2}{2 \cos^2 \theta_W M_Z^2} \left(\frac{1}{2} Z_{uc} \right)^2 (\bar{u}_L \gamma^\mu c_L) (\bar{u}_L \gamma_\mu c_L). \quad (3.1)$$

From the relations:

$$\begin{aligned} \Delta M_D &\simeq 2 |M_{12}^D| \\ (M_{12}^D)_{ZFCNC} &= -\langle D^0 | \mathcal{L}_{\text{eff}}^Z | \bar{D}^0 \rangle \end{aligned} \quad (3.2)$$

$$\langle D^0 | (\bar{u}_L \gamma^\mu c_L) (\bar{u}_L \gamma_\mu c_L) | \bar{D}^0 \rangle = \frac{1}{4} \frac{8}{3} \frac{F_D^2 M_D^2}{2 M_D} B_D \eta$$

one then obtains

$$(M_{12}^D)_{ZFCNC} = \frac{\sqrt{2} G_F F_D^2 B_D M_D}{6} Z_{uc}^2 \eta \quad (3.3)$$

where η is a QCD correction factor expected to be of order one.

The size of the new contribution to $D^0 - \bar{D}^0$ mixing depends crucially on the magnitude of Z_{uc} . From Eqs. (2.8), (2.5) it follows that the Z_{ij} are related to the deviations from unitarity of V_{CKM} . One can then use the experimental knowledge of V_{CKM} to constrain the size of Z_{uc} . Explicitly, one has

$$|S_1|^2 = 1 - |V_{ud}|^2 - |V_{us}|^2 - |V_{ub}|^2. \quad (3.4)$$

From the experimental values [11]:

$$\begin{aligned} |V_{ud}| &= 0.9744 \pm 0.0010 \\ |V_{us}| &= 0.2205 \pm 0.0018 \\ |V_{ub}|/|V_{cb}| &= 0.08 \pm 0.02 \\ |V_{cb}| &= 0.040 \pm 0.005 \end{aligned} \quad (3.5)$$

it follows that:

$$|S_1| \leq 7 \cdot 10^{-2}. \quad (3.6)$$

A bound on $|S_2|$ may also be derived from Eq. (2.5). However, due to the poor experimental knowledge on $|V_{cs}|$, only a loose bound on $|S_2|$ can be obtained

$$|S_2| \leq 0.5. \quad (3.7)$$

One may also try to obtain a limit on $|S_2^* S_3|$ using the relation:

$$|S_2^* S_3| = |V_{cd} V_{td}^* + V_{cs} V_{ts}^* + V_{cb} V_{tb}^*|. \quad (3.8)$$

However, given the present knowledge of V_{CKM} , only a loose bound can be obtained. If, for example, one assumes $V_{ts} \approx V_{cb}$, it follows that

$$|S_2^* S_3| < 10^{-1}. \quad (3.9)$$

Next we will analyse the expectations for the strength of the Z_{ij} , within the present model. The magnitude of the Z_{ij} crucially depends on the assumptions one makes for the J_i , defined by Eq. (2.1). In order to illustrate this dependence, we will analyse the size of the Z_{ij} , for two special cases:

i) Non-Hierarchical J_i 's

The simplest assumption one can make about the J_i is that they are all roughly of the same order of magnitude, i.e. $J_1 \sim J_2 \sim J_3 \equiv J$. Using Eqs. (2.8), (2.9) one then obtains:

$$|Z_{uc}| \simeq |Z_{ct}| \simeq |Z_{ut}| \simeq \frac{|J|^2}{M^2}. \quad (3.10)$$

In this case, the best bound on $|J|/M$ arises from the experimental bound on ΔM_D [11]:

$$(\Delta M_D)_{\text{exp}} < 1.3 \cdot 10^{-13} \text{ GeV}. \quad (3.11)$$

From Eqs. (3.1), (3.7), (3.8), one obtains

$$|J|/M < 3.3 \cdot 10^{-2} \quad (3.12)$$

where we have used $B_D F_D^2 \approx 0.01 \text{ GeV}^2$, $M_D = 1.9 \text{ GeV}$. From Eqs. (3.3), (3.7), one concludes that if one assumes, for example, $M \sim 0.5 \text{ TeV}$, then $(\Delta M_D)_{ZFCNC}$ would be two orders of magnitude larger than $(\Delta M_D)_{HQEFT}$ for $J \approx 5 \text{ GeV}$. Taking into account the size of m_t , it is clear that larger values of J are entirely plausible and therefore $(\Delta M_D)_{ZFCNC}$ can be even further enhanced.

ii) Hierarchical J_i 's

If one assumes that the J_i differ significantly from each other, the simplest possibility is to consider that the J_i follow the standard up quark hierarchy, i.e. $J_1 \sim m_u$,

$J_2 \sim m_c$, $J_3 \sim m_t$. Using Eqs. (2.9) and putting $m_c \sim 1.2$ GeV, $m_t \sim 174$ GeV, $V_{cb} \sim 0.05$, one obtains

$$|S_1| \sim \frac{1.15}{M}; \quad |S_2| \sim \frac{10}{M}; \quad |S_3| \sim \frac{174}{M}. \quad (3.13)$$

If we take $M \simeq 0.5$ TeV, we obtain:

$$Z_{uc} = |S_1 S_2| \simeq 0.46 \cdot 10^{-3}. \quad (3.14)$$

For this value of Z_{uc} , one obtains $(\Delta M_D)_{ZFCNC}$ about two orders of magnitude above the value of $(\Delta M_D)_{HQEFT}$, and therefore at the reach of the next round of experiments.

3.2 $D^0 - \bar{D}^0$ mixing in a model with a $Q = -1/3$ isosinglet quark

If the SM is extended by adding a $Q = -1/3$ isosinglet quark N , there will be ZFCNC in the down quark sector and the SM portion of the new CKM matrix will also deviate from a unitary matrix. We may choose, without loss of generality, a weak-basis where the 3×3 up quark mass matrix is diagonal, and the 4×4 down quark mass matrix \mathcal{M}_d has the same form as that given by Eq. (2.1). If we define the matrix W in this context as the unitary matrix which diagonalizes $\mathcal{M}_d \mathcal{M}_d^\dagger$, it can still be specified by Eq. (2.3) where, as before, the block S parametrizes the ZFCNC and the block R gives the new charged boson couplings. The charged current interactions are now given by:

$$\frac{g}{\sqrt{2}} [\bar{u}_L K \gamma_\mu d_L + \bar{u}_L R \gamma_\mu N] W^\mu + h.c. \quad (3.15)$$

In this case, the short distance contributions to the mixing in the $D^0 - \bar{D}^0$ system arise through the box diagrams of Fig. 2 where the quarks participating in the internal lines are the down quarks d, s, b and the new quark N , and with $R_i = V_{iN}$ in the notation of Fig. 2. The coefficients $V_{u\alpha}, V_{c\alpha}$ appearing in the vertices of Fig. 2, correspond to the first and the second row of the new (3×4) CKM matrix. Since this matrix consists of the first three lines of a 4×4 unitary matrix, the orthogonality relation $\sum V_{c\alpha}^* V_{u\alpha} = 0$ still holds and, as a result, the technical details of the calculation of these box diagrams coincide with those encountered in the framework of four standard generation model [12]. The new aspect of the present model is that the strength of the couplings of the extra down-quark is related to its mass.

In the $D^0 - \bar{D}^0$ system, the usual zero external momentum approximation in the evaluation of the box diagram [13] for three generations in the SM is not justified because the masses of the internal quarks are not always large compared to the c -quark mass and a more detailed calculation has been performed by several authors [10]. Yet, for the quark N , which we assume to be much heavier than the c -quark, the zero external momentum approximation is good and we have (neglecting CP violating effects)

$$\Delta M_D(N, N) \simeq 2|M_{12}(N, N)| \simeq \frac{G_F^2}{6\pi^2} F_D^2 B_D m_D M_W^2 |\lambda_N^2| S(x_N) \eta \quad (3.16)$$

where $\lambda_N = V_{uN}V_{cN}^*$, $x_N = (m_N/m_W)^2$, η denotes a QCD correction coefficient and $S(x_N)$ is an Inami–Lim function:

$$S(x_N) = x_N \left[\frac{1}{4} + \frac{9}{4} \frac{1}{(1-x_N)} - \frac{3}{2} \frac{1}{(1-x_N)^2} \right] + \frac{3}{2} \left[\frac{x_N}{x_N-1} \right]^3 \log x_N \quad (3.17)$$

computed for an arbitrary value of the internal quark mass. We are interested in the case where the contribution of the new quark N is dominant and therefore we neglect the contributions of the standard down-type quarks. Comparing Eq. (3.15) to Eq. (2.4) one readily concludes that the unitarity constraints on S in the model with one $Q = 2/3$ isosinglet quark coincide with the constraints on R in the model with a $Q = -1/3$ isosinglet quark. In particular, the following bounds have to be satisfied:

$$|R_1| \equiv |V_{uN}| \leq 7 \cdot 10^{-2}, \quad |R_2| \equiv |V_{cN}| \leq 0.5. \quad (3.18)$$

In leading order one has:

$$R_i \cong J_i/M. \quad (3.19)$$

In the estimation of the magnitude of the factor $|\lambda_N|^2 S(x_N)$, one has to keep in mind the following points. On the one hand, given the value of m_b , a value of m_N of order e.g. 200 GeV, would still be compatible with the approximation used in the diagonalization of $\mathcal{M}_d \mathcal{M}_d^\dagger$, since one would have $m_b^2/m_N^2 \ll 1$. On the other hand, one has to take into account the experimental constraints on ZFCNC, which in the present model, with a $Q = -1/3$ isosinglet quark, appear in the down-quark sector. The most stringent constraints arise from $K_L \rightarrow \mu^+ \mu^-$ and the size of the CP violating parameter ε , which lead to the following bounds [14]:

$$|\text{Re}(Z_{ds})| \leq 2.6 \cdot 10^{-5}, \quad |\text{Im}(Z_{ds})^2| \leq 9.2 \cdot 10^{-10}. \quad (3.20)$$

We should keep in mind that in the model with a $Q = -1/3$ isosinglet quark, one has relations entirely analogous to Eqs. (2.6), (2.8). In particular,

$$Z_{ds} = -S_1^* S_2 \quad (3.21)$$

with the S_i given by Eq. (2.6). If we assume $J_1 \sim J_2 \sim J_3 \equiv J$, the bounds of Eq. (3.20) lead to:

$$J/M < 5 \cdot 10^{-3} \quad (3.22)$$

which in turn implies:

$$\lambda_N \lesssim 2.5 \cdot 10^{-5}. \quad (3.23)$$

For this value of λ_N , the contribution of the isosinglet quark to ΔM_D would be negligible. At this point, one may ask whether it is possible at all to obtain a significant contribution to ΔM_D in models with a $Q = -1/3$ isosinglet quark. From Eqs. (3.19), (3.20), (3.21) it is clear that obtaining a sizable contribution to ΔM_D while conforming to the bounds of Eq. (3.20) requires a strong cancellation among the various terms contributing to S_1 .

Assuming that this cancellation occurs, then for $J_1 \cong 3$ GeV, $J_2 \cong 15$ GeV, $M_N \cong 200$ GeV, one obtains:

$$\lambda_N \cong 1.1 \cdot 10^{-3} \quad (3.24)$$

which leads to:

$$\Delta M_D(N, N) \sim 10^{-15} \text{ GeV} \quad (3.25)$$

which is roughly two orders of magnitude larger than $(\Delta M_D)_{HQEFT}$. However, it should be emphasized that in the model with a $Q = -1/3$ isosinglet quark, a large value for $\Delta M_D(N, N)$, consistent with the bound on Z_{ds} , can only be obtained by assuming a contrived choice of parameters leading to the suppression of S_1 . This is to be contrasted with the situation one encountered in the model with a $Q = 2/3$ isosinglet quark, where a large contribution to ΔM_D is obtained without assuming any fine-tuning of parameters.

4 Single Top Production at LEP II

In the SM, the cross section for single top quark production through the reaction $e^+e^- \rightarrow t\bar{b}e\bar{\nu}_e$ is too small [15] for the detection of single top quark to be feasible at LEP II energy and planned luminosity. In the presence of ZFCNC in the $Q = 2/3$ quark sector, the lowest order single top quark production process $e^+e^- \rightarrow t\bar{c}(t\bar{u})$ is possible, thus providing a potential probe of ZFCNC involving the top quark. The possibility of searching for t -flavour violating neutral currents via single top quark production at e^+e^- colliders was raised about ten years ago by Hikasa [16]. However, the top quark mass is now known to be much larger than what was generally expected at that time. The overall strength and the structure of the t -flavour changing neutral current interactions can be parametrized in the following way:

$$\mathcal{L} = i \frac{g}{2 \cos \theta_W} z_{tq} \bar{q} \gamma_\mu (V_q - A_q \gamma_5) t Z^\mu + h.c. \quad (4.1)$$

where $q = u, c$. The cross section for the Z mediated single top production process is then given by:

$$\sigma_Z(e^+e^- \rightarrow t\bar{q}) = \frac{G_F^2}{8\pi} (1 - 4 \sin^2 \theta_W + 8 \sin^4 \theta_W) \frac{M_Z^4 (s - m_t^2)^2 (2s + m_t^2)}{s^2 (s - M_Z^2)^2} (V_q^2 + A_q^2) |z_{tq}|^2 \quad (4.2)$$

where, in view of the very high value of the top quark mass, we have safely neglected the Z -width and $m_q \equiv m_c, m_u$. The result of Eq. (4.2) is model independent. We can particularize to the present model where ZFCNC arise as a result of the mixing of a $Q = 2/3$ isosinglet quark with the standard quarks, by setting $A_q = V_q = 1/2$. Having in mind LEP II, we have plotted in Fig. 3, $y = \sigma/|z_{tq}|^2$ against m_t , for $\sqrt{s} = 175, 190$ GeV. If we consider an integrated luminosity of 500 pb^{-1} , then the values of $|z_{tq}|$ corresponding to five events are shown in Fig. 4. It is clear that only for relatively large values of Z_{tq} ($Z_{tq} \sim 0.1$) is the detection of single top at LEP II possible. We turn now to the question whether such a value of Z_{tq} is possible in the framework of the model

with a $Q = 2/3$ isosinglet quark. From Eqs. (2.8), (2.9) it is clear that one may have $Z_{tc} \sim (0.1)$, by choosing, e.g. $J_2 \approx J_3 \approx m_t$ and $m_Q \approx 0.5$ TeV. However, taking into account $Z_{uc} = |S_1 S_2|$, it is clear from Eqs. (2.9) that these value of J_2, J_3 would lead to a value of ΔM_D which would violate the experimental bound of Eq. (3.11), unless there is a cancellation in the various terms contributing to S_1 . Therefore, in the model with a $Q = 2/3$ isosinglet quark, only for a contrived choice of parameters can one obtain Z_{tc} sufficiently large to lead to single top production at LEP II and at the same time have Z_{uc} sufficiently small to conform to the experimental bound on ΔM_D .

5 Conclusions

We have investigated the impact of ZFCNC on $D^0 - \bar{D}^0$ mixing and on single top production, in the framework of an extension of the SM where isosinglet quarks are introduced. It was shown that in models with a $Q = 2/3$ isosinglet quark, with a mass of the order of one TeV, the new contribution to ΔM_D arising from ZFCNC can be more than two orders of magnitude larger than the SM contribution evaluated within the framework of HQEFT, and therefore at the reach of the next round of experiments. This enhancement of $D^0 - \bar{D}^0$ mixing is obtained for generic values of the parameters and does not involve any fine-tuning. On the contrary, in models with a $Q = -1/3$ isosinglet quark, due to the constraints arising from rare kaon decays and the value of ε , only for special choices of parameters leading to the necessary cancellations can one still obtain a sizable contribution to $D^0 - \bar{D}^0$ mixing. At this point, it is worth recalling that the study of CP asymmetries in B^0 decays provides the best way to test models with a $Q = -1/3$ isosinglet quark. Indeed, it has been shown [17] that even if the contribution of ZFCNC to $B^0 - \bar{B}^0$ mixing is only at the level of 20 per cent, the predictions for CP asymmetries in B^0 decays may differ drastically from those of the SM.

As far as single top production is concerned, we have seen that there is a new contribution to single-top production in e^+e^- reactions, due to t -flavour violating neutral currents, which arise in models with a $Q = 2/3$ isosinglet quark. However, it was shown that only for a somewhat contrived choice of parameters, is the strength of ZFCNC involving the top quark sufficiently large to lead to the detection of single top at LEP II energy and luminosity.

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Figure Captions

- Fig. 1** Tree diagram contribution to the $\Delta C = 2$ transition in a model with a $Q = 2/3$ isosinglet quark.
- Fig. 2** One of the box diagrams that induce a $\Delta C = 2$ transition in a model with a $Q = -1/3$ isosinglet quark.
- Fig. 3** The value of $y = \sigma/|z_{tq}|^2$ (in pb) as a function of the top quark mass for C.M. energies of 175 GeV (lower curve) and 190 GeV (upper curve).
- Fig. 4** The value of $|z_{tq}|_5$ corresponding to 5 events at the predicted LEP II integrated Luminosity of 500 pb^{-1} , as a function of the top quark mass, for C.M. energies at 175 GeV (upper curve) and 190 GeV (lower curve).

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